Effective field theory approach to $b \to s\ell\ell^{(\prime)}, \ B \to K^{(*)}\nu\bar{\nu}$ and $B \to D^{(*)}\tau\nu$ with third generation couplings

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LHCb reported anomalies in $B \to K^* \mu^+ \mu^-$, $B_s \to \phi \mu^+ \mu^-$ and $R(K) = B \to K \mu^+ \mu^-/B \to K e^+ e^-$. Furthermore, BaBar, BELLE and LHCb found hints for the violation of lepton flavour universality violation in $R(D^{(*)}) = B \to D^{(*)} \tau \nu/B \to D^{(*)} \ell \nu$. In this note we reexamine these decays and their correlations to $B \to K^{(*)} \nu \bar{\nu}$ using gauge invariant dim-6 operators. For the numerical analysis we focus on scenarios in which new physics couples, in the interaction eigenbasis, to third generation quarks and lepton only. We conclude that such a setup can explain the $b \to s \mu^+ \mu^-$ data simultaneously with $R(D^{(*)})$ for small mixing angles in the lepton sector (of the order of $\pi/16$) and very small mixing angles in the quark sector (smaller than V_{cb}). In these region of parameter space $B \to K^{(*)} \tau \mu$ and $B_s \to \tau \mu$ can be order 10^{-6} . Possible UV completions are briefly discussed.

I. INTRODUCTION

So far, the LHC completed the standard model (SM) of particle physics by discovering the last missing piece, the Higgs particle [1, 2].[66] Furthermore, no significant direct evidence for physics beyond the SM has been found, i.e. no new particles were discovered. However, LHCb observed indirect 'hints' for new physics (NP) in $B \to K^*\mu^+\mu^-$, $B_s \to \phi\mu^+\mu^-$ and $R(K) \equiv \text{Br}(B \to K\mu^+\mu^-)/\text{Br}(B \to Ke^+e^-)$. Furthermore, BaBar and also very recently BELLE and LHCb reported lepton flavour universality violation in $B \to D^{(*)}\tau\nu$. These observations can be used as a guideline in the exploration of possible physics beyond the SM.

In more detail, the current experimental situation is as follows: LHCb reported deviations from the SM predictions [3] in $B \to K^* \mu^+ \mu^-$ [4, 5] (mainly in an angular observable called P_5' [6]) with a significance of 2–3 σ depending on the assumptions of hadronic uncertainties [7–9]. Also in the decay $B_s \to \phi \mu^+ \mu^-$ [10] LHCb uncovered differences compared to the SM prediction based on lattice QCD [11, 12] and light-cone sumrules [13] of 3.1 σ [8].[67] Furthermore, LHCb [14] found indications for the violation of lepton flavour universality, namely

$$R(K) = 0.745^{+0.090}_{-0.074} \pm 0.036,$$
 (1)

in the range $1\,\mathrm{GeV^2} < q^2 < 6\,\mathrm{GeV^2}$. This measurement is in tension with the theoretically clean SM prediction $R_\mathrm{SM}(K) = 1.0003 \pm 0.0001$ [15] by $2.6\,\sigma$. Combining these anomalies with all other observables for $b \to s\mu^+\mu^-$ transitions, it is found that a scenario with NP in $C_9^{\mu\mu}$ (corresponding to the operator $\bar{s}\gamma^\nu P_L b\,\bar{\mu}\gamma_\nu\mu$) but not in C_9^{ee} is preferred compared to the SM by $4.3\,\sigma$ [16].

Hints for lepton flavour universality violating NP also comes from the BaBar collaboration that performed an analysis of the semileptonic B decays $B \to D^{(*)} \tau \nu$ [17]. Recently, these decays have also been reanalyzed by

BELLE [18] and LHCb measured $B \to D^*\tau\nu$ [19]. In summary, these experiments have found for the ratios $R(D^{(*)}) \equiv \text{Br}(B \to D^{(*)}\tau\nu)/\text{Br}(B \to D^{(*)}\ell\nu)$:

$$R(D)_{\text{BaBar}} = 0.440 \pm 0.058 \pm 0.042,$$
 (2)

$$R(D)_{\text{BELLE}} = 0.375^{+0.064}_{-0.063} \pm 0.026,$$
 (3)

$$R(D^*)_{\text{BaBar}} = 0.332 \pm 0.024 \pm 0.018,$$
 (4)

$$R(D^*)_{\text{BELLE}} = 0.293^{+0.039}_{-0.037} \pm 0.015,$$
 (5)

$$R(D^*)_{\text{LHCb}} = 0.336 \pm 0.027 \pm 0.030$$
. (6)

Here the first (second) errors are statistical (systematic). Combining these measurements one finds [20]

$$R(D)_{\text{EXP}} = 0.388 \pm 0.047,$$

 $R(D^*)_{\text{EXP}} = 0.321 \pm 0.021.$ (7)

Comparing these measurements to the SM predictions [21]

$$R_{\rm SM}(D) = 0.297 \pm 0.017,$$

 $R_{\rm SM}(D^*) = 0.252 \pm 0.003,$ (8)

we see that there is a discrepancy of $1.8\,\sigma$ for R(D) and $3.3\,\sigma$ for $R(D^*)$ and the combination corresponds approximately to a $3.8\,\sigma$ deviation from the SM (compared to $3.4\,\sigma$ taking into account the BaBar results only [17]).

Numerous models have been proposed in order to explain the anomalies in $b \to s\mu^+\mu^-$ transitions (see for example Refs. [22–32] for Z' models and Refs. [33, 34] for models with leptoquarks) and the deviations from the SM predictions in tauonic B decays [35–43].

Alternatively, a model independent approach using higher dimensional operators has been employed, as in the model independent fits [6, 8, 44]. In this context, it has been argued that as R(K) violates lepton flavour universality (LFU) also lepton flavour could be violated in B decays [45] which might be linked to neutrino oscillations [46]. [68] While [45] considered the effect of operators

at the B meson scale which are invariant under electromagnetic gauge interactions only, also operators invariant under the full SM gauge group [47, 48] have been considered in Ref. [49–52].[69] Here it has been claimed than an simultaneous explanation of R(K), R(D) and $R(D^*)$ using gauge invariant operators with left-handed fermions is possible [50, 52]. For this purpose, it was assumed that in the interaction eigenbasis only couplings to the third generation exist [45, 50] (or are enhanced by m_τ^2/m_μ^2 compared to the second one [52]), while all other couplings are generated by the misalignment between the mass and the interaction basis (or are suppressed by small lepton mass ratios [52]).

In this article we reconsider the possibility of explaining $B \to D^{(*)} \tau \nu$ and the $b \to s \mu^+ \mu^-$ data with higher dimensional gauge invariant operators, taking into account the constraints from $B \to K^{(*)} \nu \bar{\nu}$ and using the results of the global fit to $b \to s \mu \mu$ transitions. We extend the analysis of Ref. [52] and consider the possibility of lepton flavour violation (LFV) and compared to Ref. [45] we include the correlations due to $SU(2)_L$ gauge invariance and give quantitative predictions for $B \to K^{(*)} \tau \mu$ and $B_s \to \tau \mu$.

The outline is as follows: In the next section we collect the necessary formulae for the flavour observables. Sec. III discusses the gauge invariant higher dimensional operators relevant for our analysis and Sec. IV presents our numerical results. Sec. V briefly reviews some possible UV completions. Finally we conclude in Sec. VI.

II. FLAVOUR OBSERVABLES

A.
$$b \rightarrow s\mu^+\mu^-$$
 transitions

 $b \to s \ell_i \ell_j$ transitions are defined via the effective Hamiltonian

$$H_{\text{eff}}^{\ell_i \ell_j} = -\frac{4G_F}{\sqrt{2}} V_{tb} V_{ts}^* \sum_{a=9,10} \left(C_a^{\ell_i \ell_j} O_a^{\ell_i \ell_j} + C_a^{\prime} \ell_i \ell_j O_a^{\prime} \ell_i \ell_j \right) ,$$

$$O_{9(10)}^{\ell_i \ell_j} = \frac{\alpha}{4\pi} [\bar{s} \gamma^{\mu} P_L b] [\bar{\ell}_i \gamma_{\mu} (\gamma^5) \ell_j], \qquad (9)$$

where the primed operators are obtained by exchanging $L \leftrightarrow R$.

Concerning $B \to K^* \mu^+ \mu^-$, $B_s \to \phi \mu^+ \mu^-$ and $B \to K \mu^+ \mu^- / B \to K e^+ e^-$, as already noted in Ref. [22, 53], $C_9^{\mu\mu} < 0$ and $C_9^{\prime\mu\mu} = 0$ is preferred by data. However, also the possibility $C_9^{\mu\mu} = -C_{10}^{\mu\mu} < 0$ gives a good fit to data. Using the global fit of Ref. [8, 16] we see that at $(1 \sigma) \ 2 \sigma$ level

$$\begin{array}{ll} -0.53(-0.81) \geq & C_9^{\mu\mu} & \geq (-1.32) - 1.54(10) \\ -0.18(-0.35) \geq & C_9^{\mu\mu} = -C_{10}^{\mu\mu} & \geq (-0.71) - 0.91(11) \end{array}$$

Interestingly, the values of $C_9^{\mu\mu}$, $C_{10}^{\mu\mu}$ favoured by R(K) and $B \to K^*\mu^+\mu^-$ lie approximately in the same range.[70] Furthermore, a good fit to the current data does not require $C_9'^{\mu\mu}$, hence in the following we neglect operators with right-handed quark currents for simplicity.

B.
$$B \to K^{(*)} \nu \bar{\nu}$$

Following Ref. [51] we write the relevant effective Hamiltonian as

$$H_{\text{eff}}^{\nu_i \nu_j} = -\frac{4G_F}{\sqrt{2}} V_{tb} V_{ts}^* \left(C_L^{ij} O_L^{ij} + C_R^{ij} O_R^{ij} \right) \tag{12}$$

$$O_{L,R}^{ij} = \frac{\alpha}{4\pi} [\bar{s}\gamma^{\mu} P_{L,R} b] [\bar{\nu}_i \gamma_{\mu} (1 - \gamma^5) \nu_j], \qquad (13)$$

and $C_L^{\rm SM} \approx -1.47/s_w^2$. In the limit of vanishing right-handed sb current, the branching ratios normalized by the SM predictions read

$$R_{K^{(*)}}^{\nu\bar{\nu}} = \frac{1}{3} \sum_{i,j=1}^{3} \frac{\left| C_L^{ij} \right|^2}{\left| C_L^{\text{SM}} \right|^2}.$$
 (14)

The current experimental limits are $R_K^{\nu\bar{\nu}} < 4.3$ [54] and $R_{K^*}^{\nu\bar{\nu}} < 4.4$ [55].

C.
$$B \to D^{(*)} \tau \nu$$

The effective Hamiltonian for semileptonic $b \to c$ transitions is

$$H_{\text{eff}} = \frac{4G_F}{\sqrt{2}} V_{cb} C_{L\,ij}^{cb} [\bar{c}\gamma^{\mu} P_L b] [\bar{\ell}_i \gamma_{\mu} P_L \nu_j], \qquad (15)$$

with $C_{L\,ij}^{cb\,\mathrm{SM}}=\delta_{ij}$ (for massless neutrinos) taking into account only left handed vector currents. In this case the ratios of branching ratios are

$$\frac{R(D^{(*)})_{\text{EXP}}}{R(D^{(*)})_{\text{SM}}} = \frac{\sum_{j=1}^{3} \left| C_{L \, 3j}^{cb} \right|^{2}}{\sum_{j=1}^{3} \left| C_{L \, \ell j}^{cb} \right|^{2}}, \tag{16}$$

with $\ell=e,\mu$ which has to be compared to Eq. (8) and Eq. (7).

D. Lepton-flavour violating B decays

Here we give formulas for the branching ratios of LFV B decays following the analysis of Ref. [56]. We take into account only contributions from the operators $O_9^{(\prime)\ell\ell'}$ and $O_{10}^{(\prime)\ell\ell'}$ while neglecting contributions from operators with scalar currents not relevant for our analysis. For $B_s \to \ell^+\ell'^-$ (with $\ell \neq \ell'$) we use the results of Ref. [57] neglecting the mass of the lighter lepton. The branching ratios for $B \to K^{(*)}\tau^\pm\mu^\mp$, $B \to K^{(*)}\mu^\pm e^\mp$ are computed using form-factors obtained from lattice QCD in Ref. [58] (see also Refs. [12, 59]). The final results read

$$\operatorname{Br}\left[B_{s} \to \ell^{+}\ell'^{-}\right] = \frac{\tau_{B_{s}}m_{\ell}^{2}M_{B_{s}}f_{B_{s}}^{2}}{32\pi^{3}}\alpha^{2}G_{F}^{2}\left|V_{tb}V_{ts}^{*}\right|^{2}\left(1 - \frac{\operatorname{Max}[m_{\ell}^{2}, m_{\ell'}^{2}]}{M_{B_{s}}^{2}}\right)^{2}\left(\left|C_{9}^{\ell\ell'} - C_{9}^{\prime\ell\ell'}\right|^{2} + \left|C_{10}^{\ell\ell'} - C_{10}^{\prime\ell\ell'}\right|^{2}\right),$$

$$\operatorname{Br}\left[B \to K^{(*)}\ell^{+}\ell'^{-}\right] = 10^{-9}\left(a_{K^{(*)}\ell\ell'}\left|C_{9}^{\ell\ell'} + C_{9}^{\prime\ell\ell'}\right|^{2} + b_{K^{(*)}\ell\ell'}\left|C_{10}^{\ell\ell'} + C_{10}^{\prime\ell\ell'}\right|^{2}\right)$$

$$+ c_{K^{*}\ell\ell'}\left|C_{9}^{\ell\ell'} - C_{9}^{\prime\ell\ell'}\right|^{2} + d_{K^{*}\ell\ell'}\left|C_{10}^{\ell\ell'} - C_{10}^{\prime\ell\ell'}\right|^{2}\right), \tag{17}$$

with

$\ell\ell'$	$a_{K\ell\ell'}$	$b_{K\ell\ell'}$	$a_{K^*\ell\ell'}$	$b_{K^*\ell\ell'}$	$c_{K^*\ell\ell'}$	$d_{K^*\ell\ell'}$
$\tau \mu, \tau e$	9.6 ± 1.0	10.0 ± 1.3	3.0 ± 0.8	2.7 ± 0.7	16.4 ± 2.1	15.4 ± 1.9
μe	15.4 ± 3.1	15.7 ± 3.1	5.6 ± 1.9	5.6 ± 1.9	29.1 ± 4.9	29.1 ± 4.9

The formula for the branching ratio of $B_s \to \ell^+ \ell'^-$ is symmetric with respect to the exchange of $C_9^{(\prime)\ell\ell'} \leftrightarrow C_{10}^{(\prime)\ell\ell'}$, while in the case of $B \to K^{(*)}\ell^+\ell'^-$ this symmetry is broken by lepton-mass effects. There is a small difference between the theoretical prediction for the charged mode $B^+ \to K^{(*)+}\ell^+\ell'^-$ and the neutral one $B^0 \to K^{(*)0}\ell^+\ell'^-$ due to the different B-meson lifetime τ_B which we neglected fixing the numerical value of τ_B to the one of the neutral meson. Note that the results above are given for $\ell^-\ell'^+$ final states and not for the sum $\ell^\pm\ell'^\mp = \ell^-\ell'^+ + \ell^+\ell'^-$ to which the experimental constraints apply [60]. The only channel with $\tau\mu$ final states for which an experimental upper limit exists is

Br
$$\left[B^+ \to K^+ \tau^{\pm} \mu^{\mp} \right]_{\text{exp}} \le 4.8 \times 10^{-5} \,.$$
 (18)

III. GAUGE INVARIANT OPERATORS

As we have previously seen, a scenario with left-handed currents only gives a good fit to data, cf. Eq. (11). In such a scenario $SU(2)_L$ relations are necessarily present. These relations are automatically taken into account once gauge invariant operators are considered. Therefore, let us focus on 4-fermion operators with left-handed quarks and leptons. There are two such 4-fermion operators in the effective Lagrangian

$$\mathcal{L}_{\text{dim6}} = \frac{1}{\Lambda^2} \sum O_X C_X \,, \tag{19}$$

where Λ is the scale of NP, which can contribute to $b \rightarrow s\ell\ell$ transitions at tree-level [47, 48]:

$$Q_{\ell q}^{(1)} = (\bar{L}\gamma^{\mu}L) (\bar{Q}\gamma_{\mu}Q) , \quad Q_{\ell q}^{(3)} = (\bar{L}\gamma^{\mu}\tau_{I}L) (\bar{Q}\gamma_{\mu}\tau^{I}Q) ,$$
(20)

where L is the lepton doublet and Q the quark doublet and the flavour indices are not explicitly shown here. Writing these operators in terms of their $SU(2)_L$ components (i.e. up-quarks, down-quarks, charged leptons and

neutrinos) we find for the terms relevant for the processes discussed in the last section (before EW symmetry breaking)

$$\mathcal{L} \supset \frac{C_{ijkl}^{(1)}}{\Lambda^2} \left(\bar{\ell}_i \gamma^{\mu} P_L \ell_j \bar{d}_k \gamma_{\mu} P_L d_l + \bar{\nu}_i \gamma^{\mu} P_L \nu_j \bar{d}_k \gamma_{\mu} P_L d_l \right) +$$

$$\frac{C_{ijkl}^{(3)}}{\Lambda^2} \left(2 \bar{\ell}_i \gamma^{\mu} P_L \nu_j \bar{u}_k \gamma_{\mu} P_L d_l - \bar{\nu}_i \gamma^{\mu} P_L \nu_j \bar{d}_k \gamma_{\mu} P_L d_l + \bar{\ell}_i \gamma^{\mu} P_L \ell_j \bar{d}_k \gamma_{\mu} P_L d_l \right) ,$$

$$(21)$$

where $C_{ijkl}^{(1,3)}$ are the dimensionless coefficients of the operators of Eq. (20). After EW symmetry breaking the following redefinitions of the fields are performed in order to render the mass matrices diagonal

$$d_L \to D^{\dagger} d_L, \ u_L \to U^{\dagger} u_L, \ \ell_L \to L^{\dagger} \ell_L, \ \nu \to L^{\dagger} \nu$$
. (22)

We define for future convenience

$$\lambda^{(1,3)} \tilde{X}_{ij}^{(1,3)} \tilde{Y}_{kl}^{(1,3)} = L_{i'i}^* L_{j'j} D_{k'k}^* D_{l'l} C_{i'j'k'l'}^{(1,3)} , \qquad (23)$$

where $\lambda^{(1,3)}$ are overall constants. Using constraints from the measured CKM matrix, i.e. $V=U^{\dagger}D$, we finally obtain

$$C_{9}^{ij} = -C_{10}^{ij}$$

$$= \frac{\pi}{\sqrt{2}\Lambda^{2}G_{F}\alpha V_{tb}V_{ts}^{*}} \left(\lambda^{(1)}\tilde{X}_{ij}^{(1)}\tilde{Y}_{23}^{(1)} + \lambda^{(3)}\tilde{X}_{ij}^{(3)}\tilde{Y}_{23}^{(3)}\right)$$

$$C_{L}^{ij} = \frac{\pi}{\sqrt{2}\Lambda^{2}G_{F}\alpha V_{tb}V_{ts}^{*}} \left(\lambda^{(1)}\tilde{X}_{ij}^{(1)}\tilde{Y}_{23}^{(1)} - \lambda^{(3)}\tilde{X}_{ij}^{(3)}\tilde{Y}_{23}^{(3)}\right),$$

$$C_{L\ ij}^{cb} = -\frac{\lambda^{(3)}}{\sqrt{2}\Lambda^{2}G_{F}}\frac{\tilde{X}_{ij}^{(3)}}{V_{cb}}\sum_{I}\left(V_{2k}\tilde{Y}_{k3}^{(3)}\right),$$
(24)

for the Wilson coefficients relevant for $b\to s\mu^+\mu^-$, $B\to K^{(*)}\nu\bar{\nu}$ and $B\to D^{(*)}\tau\nu$ respectively. Note that in the limit $C^{(1)}=C^{(3)}$ the contribution to $B\to K^{(*)}\nu\bar{\nu}$ vanishes.

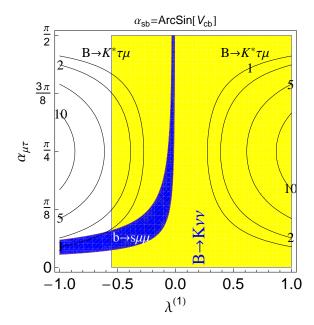


FIG. 1: Allowed regions in the $\lambda^{(1)}$ - $\alpha_{\mu\tau}$ plane from $b \to s\mu^+\mu^-$ data (blue) and $B \to K\nu\bar{\nu}$ (yellow) for $\alpha_{sb} = \text{ArcSin}[V_{cb}]$ and $\Lambda = 1 \text{ TeV}$. Note that here changing α_{sb} only has the effect of an overall scaling of $\lambda^{(1)}$. The contour lines denote $\text{Br}[B \to K^*\tau\mu]$ in units of 10^{-6} .

IV. NUMERICAL ANALYSIS

Since we have $C_9^{\tau\mu}=-C_{10}^{\tau\mu}$ we find for the LFV B decays

$$\mathrm{Br}\left[B\to K\tau^\pm\mu^\mp\right]/\mathrm{Br}\left[B\to K^*\tau^\pm\mu^\mp\right]\approx 1\,, \qquad (25)$$

Br
$$\left[B_s \to \tau^{\pm} \mu^{\mp}\right] / \text{Br} \left[B \to K^* \tau^{\pm} \mu^{\mp}\right] \approx 0.5$$
. (26)

Therefore in the following, we will just present the numerical evaluation of $\operatorname{Br}[B \to K^*\tau^{\pm}\mu^{\mp}]$ while $\operatorname{Br}[B_s \to \tau^{\pm}\mu^{\mp}]$ and $\operatorname{Br}[B \to K\tau^{\pm}\mu^{\mp}]$ can be obtained by the appropriate rescaling.

We also note that $B \to K \nu \bar{\nu}$ imposes an upper limit on the absolute value of $C_9^{\tau\tau} = -C_{10}^{\tau\tau}$ and $C_9^{\tau\mu} = -C_{10}^{\tau\mu}$ valid for $C^{(3)}$ and $C^{(1)}$ separately. Neglecting the small NP contribution to $C_L^{\mu\mu}$ and assuming no NP in the electron channel we find:

$$\frac{|C_9^{\tau\mu}|}{C_I^{\rm SM}} \le \sqrt{4.3 \times 3/2} \approx 2.5 \,, \tag{27}$$

$$\frac{|C_9^{\tau\tau}|}{C_L^{\rm SM}} \le \sqrt{3 \times 4.3 \times 3/2 - 2} + 1 \approx 5.2. \tag{28}$$

This leads to the following upper limits valid in any model generating only $C^{(3)}$ or $C^{(1)}$:

$$Br[B \to K\tau\mu] \le 8.3 \times 10^{-6}$$
. (29)

However, this limit can be evaded for $C^{(3)}=C^{(1)}$. In Ref. [52] it was proposed that the MFV-like relation $\tilde{Y}_{22}/\tilde{Y}_{33}=m_{\tau}^2/m_{\mu}^2$ could explain $R(D^{(*)})$ and $b\to$

 $s\mu^+\mu^-$ data simultaneously. From Eq. (27) we see that this ansatz is only possible for $C^{(3)} = C^{(1)}$ but not if $C^{(3)}$ or $C^{(1)}$ are separately different from zero.

Therefore, we will focus in the following on scenarios with third generation couplings in the EW basis only, which correspond to a general rank 1 matrix in the mass eigenbasis, as suggested in Ref. [45, 50]. In other words we have

$$C_{ijkl}^{(1,3)} = \lambda^{(1,3)} \tilde{X}_{ij} \tilde{Y}_{kl} ,$$
 (30)

$$\tilde{X} = L^{\dagger} X L, \ \tilde{Y} = D^{\dagger} Y D, \quad X = Y = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 1 \end{pmatrix}.$$

Taking into account only rotations among the second and third generation one finds

$$\tilde{X} = \begin{pmatrix} 0 & 0 & 0 \\ 0 & \sin^2\left(\alpha_{\mu\tau}\right) & -\sin\left(\alpha_{\mu\tau}\right)\cos\left(\alpha_{\mu\tau}\right) \\ 0 & -\sin\left(\alpha_{\mu\tau}\right)\cos\left(\alpha_{\mu\tau}\right) & \cos^2\left(\alpha_{\mu\tau}\right) \end{pmatrix},$$

$$\tilde{Y} = \begin{pmatrix} 0 & 0 & 0 \\ 0 & \sin^2\left(\alpha_{sb}\right) & -\sin\left(\alpha_{sb}\right)\cos\left(\alpha_{sb}\right) \\ 0 & -\sin\left(\alpha_{sb}\right)\cos\left(\alpha_{sb}\right) & \cos^2\left(\alpha_{sb}\right) \end{pmatrix}.$$

Note that a rotation $\sin(\alpha_{sb}) \gg V_{cb}$ would require finetuning with the up sector in order to obtain the correct CKM matrix.

1.
$$Q_{\ell q}^{(1)}$$
 operator

In this case we have neutral currents only. As a consequence, there is obviously no effect in $R(D^{(*)})$, but $b \to s\mu^+\mu^-$ is directly correlated to $B \to K^{(*)}\nu\bar{\nu}$ depending on the angle $\alpha_{\mu\tau}$. Note that a change in α_{sb} can be compensated by a change in $\lambda^{(1)}$ and therefore does not affect the correlations among $B \to K^{(*)}\nu\bar{\nu}$ and $b \to s\mu^+\mu^-$ transitions. In Fig. 1 the regions favoured by $b \to s\mu^+\mu^-$ (blue) and allowed by $B \to K\nu\bar{\nu}$ (yellow) are shown together with contour lines for $B \to K^*\tau\mu$ in units of 10^{-6} . Note that $B \to K\nu\bar{\nu}$ rules out branching ratios for $B \to K^*\tau\mu$ above approximately 1×10^{-6} and that the constraint from $B \to K\nu\bar{\nu}$, being inclusive in the neutrino flavours, is independent of $\alpha_{\mu\tau}$.

2.
$$Q_{\ell q}^{(3)}$$
 operator

Here we have also charged currents that are related to the neutral current processes via CKM rotations. In Fig. 2 the regions allowed by $B \to K \nu \bar{\nu}$ (yellow) and giving a good fit to data for $b \to s \mu^+ \mu^-$ (blue) and (at the 2σ level) for $B \to D^* \tau \nu$ (red) are shown for different values of $\lambda^{(3)}$. Note that $b \to s \mu^+ \mu^-$ data can be explained simultaneously with $R(D^{(*)})$ for negative $\mathcal{O}(1)$ values of $\lambda^{(3)}$ without violating the bounds from $B \to K \nu \bar{\nu}$. Again, in the regions compatible with all experimental constraints, the branching rations of LFV B decays to $\tau \mu$ final states can only be up to $\approx 10^{-6}$.

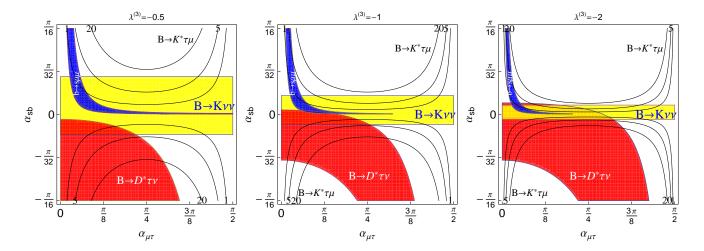


FIG. 2: Allowed regions in the $\alpha_{\mu\tau}$ - α_{sb} plane from $B \to K\nu\bar{\nu}$ (yellow), $R(D^*)$ (red) and $b \to s\mu^+\mu^-$ (blue) for $\Lambda = 1 \text{ TeV}$ and $\lambda^{(3)} = -0.5$ (left plot), $\lambda^{(3)} = -1$ (middle) and $\lambda^{(3)} = -2$ (right). Note that $\alpha_{sb} = \pi/64$ roughly corresponds to the angle needed to generate V_{cb} and that if $\lambda^{(3)}$ is positive, $R(D^*)$ and $b \to s\mu^+\mu^-$ cannot be explained simultaneously.

3.
$$Q_{\ell q}^{(1)}$$
 and $Q_{\ell q}^{(3)}$ with $\lambda^{(1)} = \lambda^{(3)}$

In this case the phenomenology is then rather similar to the case of $C^{(3)}$ only. The major differences are that, as already mentioned before, the bounds from $B \to K \nu \bar{\nu}$ are evaded and the relative contribution to $b \to s \mu \mu$ compared to $R(D^{(*)})$ is a factor of 2 larger. In Fig. 3 we show the analogous plot to the central panel of Fig. 2 $(\lambda^{(3)} = \lambda^{(1)} = -1)$ for this scenario. Note that again $R(D^{(*)})$ rules out very large branching ratios for lepton flavour violating B decays in the regions compatible with $b \to s \mu^+ \mu^-$ data. We also consider the MFV-like ansatz [52] with additional flavour rotations (light blue) which however differs only slightly for the ansatz with third generation couplings.

V. UV COMPLETIONS

Let us briefly discuss UV completions which can give the desired coupling structure. As discussed previously, the 4-Fermi operator $Q_{\ell q}^{(3)}$ is relevant both for R(K) and $R(D^{(*)})$. If $Q_{\ell q}^{(3)}$ is mediated by a single field, then there are only four possibilities: (i) Vector boson (VB) with the SM charges $(SU(3)_c, SU(2)_L, U(1)_Y) = (\mathbf{1}, \mathbf{3}, 0)$, (ii) Scalar leptoquark (SLQ) with $(\mathbf{3}, \mathbf{3}, -1/3)$, (iii) Vector leptoquark (VLQ) with $(\mathbf{3}, \mathbf{1}, 2/3)$, and (iv) Vector leptoquark with $(\mathbf{3}, \mathbf{3}, 2/3)$. The vector boson $(\mathbf{1}, \mathbf{3}, 0)$ induces only $Q_{\ell q}^{(3)}$. On the other hand, the leptoquark fields result in particular combinations of $Q_{\ell q}^{(1)}$ and $Q_{\ell q}^{(3)}$ [52]. With the assumption of the third generation coupling, the relative size of the effective couplings $\lambda^{(1,3)}$ and the signs are determined as

$$VB(\mathbf{1},\mathbf{3},0): \lambda^{(3)}$$
 both positive and negative, (31)

$$SLQ(3,3,-1/3): \lambda^{(1)} = 3\lambda^{(3)}, \quad \lambda^{(3)} > 0,$$
 (32)

$$VLQ(3,1,2/3): \lambda^{(1)} = \lambda^{(3)}, \quad \lambda^{(3)} < 0,$$
 (33)

$$VLQ(3,3,2/3): \lambda^{(1)} = -3\lambda^{(3)}, \quad \lambda^{(3)} > 0.$$
 (34)

The coefficient C_9^{ij} is proportional to $\lambda^{(1)} + \lambda^{(3)}$ and a negative value is favoured by R(K). Therefore, the scalar leptoquark is rejected as a candidate. To explain $R(D^{(*)})$ simultaneously, $\lambda^{(3)}$ itself must also be negative. This condition excludes the triplet vector leptoquark. If the experimental results are explained by the operator $Q_{\ell q}^{(3)}$ under the assumption of third generation coupling only, the possible mediators are the triplet vector boson or the singlet vector leptoquark. According to the analysis of the previous section, a good fit to flavour data requires a mediator mass of $\mathcal{O}(1)$ TeV. This opens interesting prospects for the LHC, especially in the case of leptoquarks that can be produced in proton-proton collisions via colour interactions and would decay to one lepton (τ or more interestingly μ) and one jet (possibly a b-jet).

VI. CONCLUSIONS

In this article we considered the effect of gauge invariant dim-6 operators with left-handed fermions on $b \to s \mu^+ \mu^-$, $B \to K^{(*)} \nu \bar{\nu}$, $B \to D^{(*)} \tau \nu$, $B \to K^{(*)} \tau \mu$ and $B_s \to \tau \mu$. For operators with left-handed quarks and leptons we find the correlations $\text{Br}\left[B \to K \tau^\pm \mu^\mp\right] \approx \text{Br}\left[B \to K^* \tau^\pm \mu^\mp\right] \approx 2 \text{Br}\left[B_s \to \tau^\pm \mu^\mp\right]$. We showed that the anomalies in $b \to s \mu \mu$ data can be explained simultaneously with $R(D^*)$. For this we considered scenarios in which third generation couplings in the EW basis are present only: $\lambda^{(1)} \neq 0$, $\lambda^{(3)} \neq 0$ and $\lambda^{(3)} = \lambda^{(1)} \neq 0$. Taking into account $\lambda^{(1)} \neq 0$ only, $b \to s \mu^+ \mu^-$ data can be explained without violating bounds from $B \to s \mu^+ \mu^-$

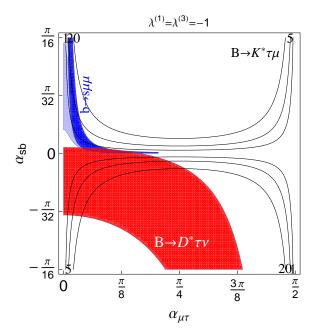


FIG. 3: Allowed regions in the $\alpha_{\mu\tau}$ - α_{sb} plane from $R(D^*)$ (red) and $b \to s\mu^+\mu^-$ (dark blue) for Λ = TeV and $\lambda^{(3)} = \lambda^{(1)} = -1$. The light blue region corresponds to the MFV-like ansatz for the lepton masses. Note that $\alpha_{sb} = \pi/64$ roughly corresponds to the angle needed to generate V_{cb} and that the MFV-like Ansatz only differs marginally from the one with third generation couplings only in the region compatible with R(D). The contour lines denote $\text{Br}[B \to K^*\tau\mu]$ in units of 10^{-6} .

 $K^{(*)}\nu\bar{\nu}$. However, in the allowed regions of parameter space, ${\rm Br}[B\to K^{(*)}\tau\mu]$ can only be up to 1×10^{-6} . In the case of $\lambda^{(3)}\neq 0,\,b\to s\mu^+\mu^-$ data can be explained simultaneously with $R(D^*)$. In these regions ${\rm Br}[B\to K^{(*)}\tau\mu]$ can again be only up to 10^{-6} . Finally we considered $\lambda^{(3)}=\lambda^{(1)}\neq 0$. Such a scenario can be realized with a leptoquark in the singlet representation of $SU(2)_L$ (making an MFV-like ansatz for the lepton couplings possible) and constraints from $B\to K^{(*)}\nu\bar{\nu}$ are avoided. Again, LFV B decays turn out to be of the same order as in the other scenarios.

Note added — During the completion of this work, an article presenting a dynamical model with additional vector bosons and third generation couplings appeared in which $Q_{\ell a}^{(3)}$ is generated [61].

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